

NONPERTURBATIVE QCD EFFECTS IN HIGH ENERGY COLLISIONS

by

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Abstract:

High energy hadron-hadron collisions are discussed. It is argued that soft collisions should involve in an essential way nonperturbative QCD. A way is outlined how to calculate properties of high energy elastic hadron-hadron scattering using field theoretic methods. The functional integrals occurring there are evaluated using the “stochastic vacuum model”. A satisfactory comparison between theory and experiment is achieved. Then the question of possible nonperturbative QCD effects in high energy hard hadron-hadron collisions is raised. It is shown that some spin effects in the Drell-Yan process may give a hint that such effects exist indeed in nature.

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1 Introduction

Today, most of us will agree that all the phenomena of hadron physics should be described in the framework of QCD. We know the simple and elegant Lagrangian density of QCD, \mathcal{L}_{QCD} , and in principle everything is derivable from it. We have

$$\mathcal{L}_{\text{QCD}}(x) = -\frac{1}{4}G_{\lambda\rho}^a(x)G^{a\lambda\rho}(x) + \sum_q \bar{q}(x)(i\gamma^\lambda D_\lambda - m_q)q(x), \quad (1.1)$$

where $G_{\lambda\rho}^a(x)$ ($a = 1, \dots, 8$) are the components of the gluon field strength tensor, $q(x)$ are the quark fields, m_q the quark masses and D_λ the covariant derivative. (All our notation follows [1]).

But the degrees of freedom in the Lagrangian density are quarks and gluons, not the hadrons we observe in nature. To make quantitative predictions for the real world starting from \mathcal{L}_{QCD} is not easy. In two areas this has been successful:

1. for short distance phenomena, where due to asymptotic freedom perturbation theory can be applied;
2. for hadron spectroscopy and other long distance phenomena, where numerical, nonperturbative methods can be applied, e.g. Monte Carlo simulations on a lattice or the discretized light cone quantization procedure.

There is a third class of phenomena which are neither pure short distance nor pure long distance: high energy hadron-hadron collisions which have contributed so much to our understanding of hadronic phenomena. These reactions are traditionally classified into “hard” and “soft” ones:

3. High energy hadron-hadron collisions:

(3.a) hard reactions,

(3.b) soft reactions.

Typical hard reactions are the Drell-Yan-type processes, e.g.

$$\begin{aligned} \pi^- + N &\rightarrow \gamma^* + X, \\ &\hookrightarrow \ell^+ \ell^- \end{aligned} \quad (1.2)$$

where $\ell = e, \mu$. All energies and momentum transfers are assumed to be large. However, the masses of the π^- and N in the initial state stay fixed and thus we are not dealing with a pure short distance phenomenon.

In the reaction (1.2) we claim to see directly the fundamental quanta of the theory, the partons, i.e. the quarks and gluons, in action (cf. Fig. 1). In the usual theoretical framework for hard reactions, the QCD improved parton model (cf. e.g. [2]), one describes the reaction of the partons, in the Drell-Yan case the $q\bar{q}$ annihilation into a virtual photon, by perturbation theory. This should be reliable, since the parton process involves only high energies and high momentum transfers. All the long distance physics due to the bound state nature of the hadrons is then lumped into parton distribution functions of the participating hadrons. This is called the factorization hypothesis, which after early investigations of soft initial and final state interactions [3] was formulated and studied in low orders of QCD perturbation theory in [4]. Subsequently, great theoretical effort has gone into proving factorization in the framework of QCD perturbation theory [5]-[7]. The result seems to be that factorization is most probably correct there. However, it is legitimate to ask if factorization is respected also by nonperturbative effects. To my knowledge this question was first asked in [8]-[10]. In Sect. 4 of this seminar I will come back to this question and will argue that there may be evidence for a breakdown of factorization in the Drell-Yan reaction.

The other class of reactions I will discuss in the following are soft high energy collisions. A typical reaction we will be interested in is proton-proton elastic scattering:

$$p + p \rightarrow p + p \quad (1.3)$$

at c.m. energies $E_{\text{cm}} = \sqrt{s} \gtrsim 5 \text{ GeV}$ say and small momentum transfers $\sqrt{|t|} = |q| \lesssim 1 \text{ GeV}$. Here we have two scales, one staying finite, one going to infinity:

$$\begin{aligned} E_{\text{cm}} &\rightarrow \infty, \\ |q| &\lesssim 1 \text{ GeV}. \end{aligned} \quad (1.4)$$

Thus, none of the above mentioned calculational methods is directly applicable. Indeed, most theoretical papers dealing with reactions in this class develop and apply models which are partly older than QCD, partly QCD “motivated”. Let me list some models for hadron-hadron elastic scattering at high energies:

- geometric [11],
- eikonal [12],
- additive quark model [13],
- Regge poles [14],
- topological expansions and strings [15],
- valons [16],
- leading log summations [17],

two-gluon exchange [18],
etc.

It would be a forbidding task to collect all references in this field. The references given above should thus only be considered as representative ones. In addition I would like to mention the inspiring general field theoretic considerations for high energy scattering and particle production by Heisenberg [19] and the impressive work by Cheng and Wu on high energy behaviour in field theories in the framework of perturbative calculations [20].

I will now argue that the theoretical description of measurable quantities of soft high energy reactions like the total cross sections should involve in an essential way nonperturbative QCD. To see this, consider massless pure gluon theory where all “hadrons” are massive glueballs. Then we know from the renormalization group analysis that the glueball masses must behave as

$$m_{\text{glueball}} \propto M e^{-c/g^2(M)} \quad (1.5)$$

for $M \rightarrow \infty$, i.e. for $g(M) \rightarrow 0$, due to asymptotic freedom. Here M is the renormalization scale, $g(M)$ is the QCD coupling strength at this scale and c is a constant. Masses in massless Yang-Mills theory are a purely nonperturbative phenomenon, due to “dimensional transmutation”. Scattering of glueball-hadrons in massless pure gluon theory should look very similar to scattering of hadrons in the real world, with finite total cross sections, amplitudes with analytic t dependence etc. At least, this would be my expectation. If the total cross section σ_{tot} has a finite limit as $s \rightarrow \infty$ we must have from the same renormalization group arguments:

$$\lim_{s \rightarrow \infty} \sigma_{\text{tot}}(s) \propto M^{-2} e^{2c/g^2(M)} \quad (1.6)$$

for $g(M) \rightarrow 0$. In this case, the total cross sections in pure gluon theory are also nonperturbative objects! It is easy to see that this conclusion is not changed if $\sigma_{\text{tot}}(s)$ has a logarithmic behaviour with s for $s \rightarrow \infty$, e.g.

$$\sigma_{\text{tot}}(s) \rightarrow \text{const} \times (\log s)^2. \quad (1.7)$$

I would then expect that also in full QCD total cross sections are nonperturbative objects, at least as far as hadrons made out of light quarks are concerned.

Some time ago P.V. Landshoff and myself started to think about a possible connection between the nontrivial vacuum structure of QCD - a typical nonperturbative phenomenon - and soft high energy reactions [21]. In the following I will first review some common folklore on the QCD vacuum and then sketch possible consequences of these ideas for high energy collisions.

2 The QCD Vacuum

According to current theoretical prejudice the vacuum state in QCD has a very complicated structure [22]-[32]. It was first noted by Savvidy [22] that by introducing a constant chromomagnetic field

$$\mathbf{B}^a = \mathbf{n}\eta^a B, \quad (a = 1, \dots, 8), \quad (2.1)$$

into the perturbative vacuum one can lower the vacuum-energy density $\varepsilon(B)$. Here \mathbf{n} and η^a are constant unit vectors in ordinary and colour space. The result of his one-loop calculation was

$$\varepsilon(B) = \frac{1}{2}B^2 + \frac{bg^2}{32\pi^2}B^2 \left[\ln \frac{B}{M^2} - \frac{1}{2} \right] \quad (2.2)$$

where g is the strong coupling constant, M is again the renormalization scale, and b is given by the lowest order term in the Callan-Symanzik β -function:

$$\beta(g) = -\frac{b}{16\pi^2}g^3 + \dots \quad (2.3)$$

For 3 colours and f flavours:

$$b = 11 - \frac{2}{3}f \quad (2.4)$$

thus, as long as we have asymptotic freedom, i.e. for $f \leq 16$, the energy density $\varepsilon(B)$ looks as indicated schematically in Fig. 2 and has its minimum for $B = B_{\text{vac}} \neq 0$. Therefore, we should expect the QCD-vacuum to develop spontaneously a chromomagnetic field, the situation being similar to that in a ferromagnet below the Curie temperature where we have spontaneous magnetization.

Of course, the vacuum state in QCD has to be relativistically invariant and cannot have a preferred direction in ordinary space and colour space. What has been considered [28] are states composed of domains with random orientation of the gluon-field strength (Fig. 3). This is analogous to Weiss domains in a ferromagnet. The vacuum state should then be a suitable linear superposition of states with various domains and orientation of the fields inside the domains. This implies that the orientation of the fields in the domains as well as the boundaries of the domains will fluctuate.

A very detailed picture for the QCD vacuum along these lines has been developed in ref. [28]. I cannot refrain from comparing this modern picture of the QCD vacuum (Fig. 4a) with the “modern picture” of the ether developed by Maxwell more than 100 years ago (Fig. 4b). The analogy is quite striking and suggests to

me that with time passing on we may also be able to find simpler views on the QCD vacuum as Einstein did with the ether. In the following we will adopt the picture of the QCD vacuum as developed in refs. [22]-[29],[31] and outlined above as a working hypothesis.

Let me now come to the values for the field strengths \mathbf{E}^a and \mathbf{B}^a in the vacuum. These must also be determined by Λ , the QCD scale parameter, the only dimensional parameter in QCD if we disregard the quark masses. Therefore, we must have on dimensional grounds for the renormalization group invariant quantity $(gB)^2$

$$(gB)^2 \sim \Lambda^4. \quad (2.5)$$

But we have much more detailed information on the values of these field strengths due to the work of Shifman, Vainshtein, and Zakharov (SVZ), who introduced the gluon condensate and first estimated its value using sum rules for charmonium states [25]:

$$\begin{aligned} \langle 0 | \frac{g^2}{4\pi^2} G_{\mu\nu}^a(x) G^{\mu\nu a}(x) | 0 \rangle &\equiv \langle 0 | \frac{g^2}{2\pi^2} (\mathbf{B}^a(x) \mathbf{B}^a(x) - \mathbf{E}^a(x) \mathbf{E}^a(x)) | 0 \rangle \\ &\equiv G_2 = (2.4 \pm 1.1) \cdot 10^{-2} \text{GeV}^4 \\ &= (335 - 430 \text{MeV})^4. \end{aligned} \quad (2.6)$$

Here we quote numerical values as given in the review [34]. A simple analysis shows that this implies

$$\langle 0 | g^2 \mathbf{B}^a(x) \mathbf{B}^a(x) | 0 \rangle = - \langle 0 | g^2 \mathbf{E}^a(x) \mathbf{E}^a(x) | 0 \rangle = \pi^2 G_2 \simeq (700 \text{MeV})^4. \quad (2.7)$$

To prove eq. (2.7) we note that Lorentz- and Parity-invariance require the vacuum expectation value of the uncontracted product of two gluon field strengths to be of the form

$$\langle 0 | \frac{g^2}{4\pi^2} G_{\mu\nu}^a(x) G_{\rho\sigma}^b(x) | 0 \rangle = (g_{\mu\rho} g_{\nu\sigma} - g_{\mu\sigma} g_{\nu\rho}) \delta^{ab} \frac{G_2}{96} \quad (2.8)$$

where G_2 is the same constant as in (2.6). Taking appropriate contractions leads to (2.6) and (2.7).

We find that $\langle 0 | \mathbf{B}^a(x) \mathbf{B}^a(x) | 0 \rangle$ is positive, $\langle 0 | \mathbf{E}^a(x) \mathbf{E}^a(x) | 0 \rangle$ negative! This can happen because we are really considering products of field operators, normal-ordered with respect to the perturbative vacuum. The interpretation of eq. (2.7) is, therefore, that the B-field fluctuates with bigger amplitude, the E-field with smaller amplitude than in the perturbative vacuum state.

What about the size a of the colour domains and the fluctuation times τ of the colour fields? On dimensional grounds we must have

$$a \sim \tau \sim \Lambda^{-1}. \quad (2.9)$$

A detailed model for the QCD vacuum incorporating the gluon condensate idea and a fall-off of the correlation of two field strengths with distance was proposed in [35]: the “stochastic vacuum model”. There the basic object is the vacuum expectation value of the “gluon field strength correlator”:

$$< 0|G_{\mu\nu}^a(x) \text{ string operator } G_{\rho\sigma}^b(y)|0 > . \quad (2.10)$$

This is the vacuum expectation value of the product of two field strengths at different points x, y made gauge-invariant by the insertion of a suitable Schwinger string operator [37], [34]. Now an ansatz for the function (2.10) is made. The fall-off of this function as $|x - y|$ gets large can be used to give a precise meaning to a (2.9) as a correlation length. Various measurable quantities can be calculated within this model, and the parameters of the above-mentioned ansatz can be determined by comparison with experiment. In this way one finds from low energy phenomenology [34]

$$a \simeq 0.35 \text{ fm}. \quad (2.11)$$

This is smaller, albeit not much smaller, than a typical radius R of a light hadron (cf. e. g. [36]):

$$R \sim 0.7 - 1 \text{ fm}. \quad (2.12)$$

Still

$$a^2/R^2 \approx 0.2 - 0.3 \quad (2.13)$$

is a reasonably small number and this will be important for us in the following.

The stochastic vacuum model is quite interesting as it leads for instance to an easy and intuitive understanding of confinement [35]. For us here it will be a tool to evaluate functional integrals in the nonperturbative domain.

3 Soft Hadronic Reactions

Consider now a quark sailing through the vacuum. The quark will interact with the gluons in the vacuum, it will have a gluon cloud around it, where the strength of the gluon cloud is governed by $g^2 < G_{\mu\nu} G^{\mu\nu} >$ and its extension in transverse directions by a . Suppose a second quark comes in opposite direction. The interaction of the two quarks happens in essence in the following way: one quark changes the gluon distribution in the vacuum, this change is felt by the other quark. Thus we estimate naively that in this picture two quarks must come closer than a distance of order a in transverse directions for scattering to occur (Fig. 5).

This quark-quark encounter could, for instance, take place in a pp -elastic scattering at high energies (Fig. 6). since $a^2 \ll R^2$ (2.13), the chance to find two quarks

of one proton simultaneously at a transverse distance $\lesssim a$ of a quark in the other proton is small. We arrive at the conclusion that single quark scattering should dominate! This would be a simple explanation of the success of the additive quark rule [13] found nearly 30 years ago. Note that it is very difficult to understand this rule in perturbative QCD where gluon exchange leads in essence to a long-range Coulomb-type force where formally $a \rightarrow \infty$.

In [21] P. V. Landshoff and myself investigated an abelian gluon model which realized the above simple physical picture. In [38] I developed these ideas further, starting from \mathcal{L}_{QCD} and trying to use only “honest” field-theoretic methods. I could show that the quark-quark scattering amplitude can be obtained by first calculating the amplitudes for each quark q_1, q_2 scattering in an external gluon potential (Fig. 7). Let the amplitudes of this scattering be $M_1(G), M_2(G)$ for the quarks q_1, q_2 , respectively. In a second step one has to average the product of these two scattering amplitudes over all gluon potentials with an appropriate functional integral measure:

$$\langle q_1 q_2 | T | q_1 q_2 \rangle \propto \langle M_1(G) M_2(G) \rangle_{G\text{-average}} . \quad (3.1)$$

Furthermore I found that in the high-energy limit $M_{1,2}(G)$ are governed by the nonabelian phase factors

$$V_{1,2}(G) \sim P \left\{ \exp -ig \int_{1,2} dz_\mu G^\mu(z) \right\} , \quad (3.2)$$

where

$$G^\mu(z) = G^{a\mu}(z) \frac{\lambda_a}{2} \quad (3.3)$$

is the gluon potential matrix, and P means path ordering. The line integrals run along the paths travelled by the quarks $q_{1,2}$ in Minkowski space. Inserting this in (3.1) leads to

$$\langle q_1 q_2 | T | q_1 q_2 \rangle \sim \langle V_1(G) V_2(G) \rangle_{G\text{-average}} . \quad (3.4)$$

How can we discuss such an expression further? Equation (3.4) means that we have to take a whole sum of gluon interactions of two quark lines and average them (Fig. 8). Surely we do not want to go back to the perturbative expansion in powers of g from there, since we remember our argument (1.6). Instead we can make a nonperturbative ansatz for the gluon propagator and in an abelian model reproduce the results of [21] in an easy way. To generalize this for full QCD with its non-abelian character we will use the methods of the stochastic vacuum model. The original version of this model was formulated by Dosch and Simonov for functional integrals in Euclidian space [35]. However, to evaluate the functional integral on the r.h.s. of (3.4) we have to work in Minkowski space — at least we found no way

to make a sort of Wick rotation in this case. To deal with such a situation, Dosch and Krämer [39] extended the stochastic vacuum model. They made an analytic continuation — i.e. a Wick rotation — for the correlators (2.10), originally defined in Euclidian space, from there to Minkowski space. They also introduced a suitable factorization scheme for higher point correlators of gluon field strengths. These ansätze give us now a method to evaluate functional integrals in Minkowski space in the nonperturbative domain. The contributions to such functional integrals can be ordered according to the number of 2-point correlators appearing.

For qq -elastic scattering with no net colour exchange the simplest contribution involves two 2-point gluon correlators and has a structure reminiscent of 2-gluon exchange (Fig. 9). Thus we expect for the amplitude

$$\begin{aligned} \langle qq|T|qq \rangle|_{elastic} &\propto \langle 0|G G G G|0 \rangle \\ &\propto \langle 0|G G|0 \rangle \langle 0|G G|0 \rangle \\ &\propto G_2^2, \end{aligned} \tag{3.5}$$

where G_2 is the gluon condensate and we have used the assumption of factorization for the 4-point correlator. Our estimate for the total qq cross section is then through the optical theorem

$$\sigma_{tot}(qq) \propto const. G_2^2 a^{10}. \tag{3.6}$$

Here we have given the correct dimension to σ_{tot} by multiplying with the appropriate power of the vacuum correlation length a , the only dimensionful quantity available in this problem apart from G_2 if we set the quark masses to zero. Through the additive quark rule we estimate also:

$$\sigma_{tot}(pp) = const. G_2^2 a^{10}. \tag{3.7}$$

The explicit calculations in the framework of the Minkowskian version of the stochastic vacuum model lead to the following results [39]-[41]. For an abelian gluon theory the above estimates (3.6), (3.7) are correct, the additive quark rule for total cross sections at high energies holds. For the non-abelian theory the additive quark rule is not obtained in this framework. It turns out that the amplitude for single quarks scattering on each other is not a sensible object. Sensible objects are the amplitudes for scattering of mesons, considered as $q\bar{q}$ wave packets, on each other. Similarly, the scattering of mesons or baryons on baryons can be treated as scattering of $q\bar{q}$ or qqq on qqq wave packets. Definite rules for writing down and evaluating these scattering amplitudes can be given. Now, in addition to the vacuum parameters G_2 and a also the radii of the hadrons — i.e. of the wave packets representing them —

enter in the results. A fit to numerical results gives for the total cross section and the slope parameter at $t = 0$ of elastic proton-proton scattering

$$\sigma_{tot}(pp) = 0.00881 \left(\frac{R_p}{a} \right)^{3.277} \cdot (3\pi^2 G_2)^2 \cdot a^{10}, \quad (3.8)$$

$$b_{pp} := \frac{d}{dt} \ln \frac{d\sigma_{el}}{dt}(pp)|_{t=0} = 1.558a^2 + 0.454R_p^2. \quad (3.9)$$

Here R_p is the proton radius and the formulae (3.8), (3.9) are valid for

$$1 \leq R_p/a \leq 3. \quad (3.10)$$

To compare (3.8), (3.9) with experimental results, we can, for instance, consider the c.m. energy $\sqrt{s} = 20$ GeV and take as input the following measured values (cf. [41]):

$$\begin{aligned} \sigma_{tot}(pp)|_{\text{Pomeron part}} &= 35 \text{ mb}, \\ b_{pp} &= 12.5 \text{ GeV}^{-2}, \\ R_p &\equiv R_{p,elm} = 0.86 \text{ fm}. \end{aligned} \quad (3.11)$$

We obtain then from (3.8) and (3.9):

$$\begin{aligned} a &= 0.31 \text{ fm}, \\ G_2 &= 6.61 \times 10^{-2} \text{ GeV}^4. \end{aligned} \quad (3.12)$$

The correlation length a comes out in surprising good agreement with the determination of this quantity from low energy phenomenology (2.11). The gluon condensate value in (3.12) is somewhat larger than the value from sum rule determinations (2.6). There are excuses for that (cf. [41]).

But perhaps we were lucky in picking out the right c.m. energy \sqrt{s} and radius for our comparison of theory and experiment. What about the s -dependence of the total cross section σ_{tot} and slope parameter b ? The vacuum parameters G_2 and a should be independent of the energy \sqrt{s} . On the other hand, it seems quite plausible to us that the effective strong interaction radii R of hadrons may depend on \sqrt{s} . Let us consider again pp (or $p\bar{p}$) elastic scattering. Once we have fixed G_2 and a from the data at $\sqrt{s} = 20$ GeV (3.8) and (3.9) give us $\sigma_{tot}(pp)$ and b_{pp} in terms of the single parameter R_p , i.e. we obtain as prediction of the model a curve in the plane b_{pp} versus $\sigma_{tot}(pp)$. This is shown in Fig. 10. It is quite remarkable that the data from $\sqrt{s} = 20$ GeV up to Tevatron energies, $\sqrt{s} = 1.8$ TeV follow this curve. For more details and further results we refer to [41].

Summarizing this section we can say that explicit calculations for high energy-elastic hadron-hadron scattering near the forward direction have been performed combining the field-theoretic methods of [38] and the Minkowski version of the stochastic vacuum model of [39]. The results are encouraging and support the idea that the vacuum structure of QCD plays an essential role in soft high-energy scattering.

4 “Synchrotron Radiation” from the Vacuum and Spin Correlations in the Drell-Yan Reaction

Let us consider for definiteness again a proton-proton collision at high c.m. energy $\sqrt{s} \gg m_p$. We look at this collision in the c.m. system and choose as z -axis the collision axis (Fig. 11). According to Feynman’s parton dogma [42] the hadrons look like jets of almost non-interacting partons, i.e. quarks and gluons. Accepting our previous views of the QCD vacuum (Sect. 2), these partons travel in a background chromomagnetic field.

What sort of new effects might we expect to occur in this situation? Consider for instance a quark-antiquark collision in a chromomagnetic field. In our picture this is very similar to an electron-positron collision in a storage ring (Fig. 12). We know that in a storage ring e^- and e^+ are deflected and emit synchrotron radiation. They also get a transverse polarization due to emission of spin-flip synchrotron radiation [43], [44]. Quite similar we can expect the quark and antiquark to be deflected by the vacuum fields. Since quarks have electric and color charge, they should then emit both photon and gluon “synchrotron radiation”. Of course, as long as we have quarks within a single, isolated proton (or antiproton) travelling through the vacuum no emission of photons can occur, and we should consider such processes as contributing to the cloud of quasi-real photons surrounding a fast-moving proton. (This is similar in spirit to the well-known Weizsäcker-Williams approximation.) But in a collision process the parent quark or antiquark will be scattered away and the photons of the cloud can become real, manifesting themselves as prompt photons in hadron-hadron collisions.

In ref. 10 we have given an estimate for the rate and the spectrum of such prompt photons using the classical formulae for synchrotron radiation [44]. The result we found can be summarized as follows: In the overall c.m. system of the hadron-hadron collision “synchrotron” photons should appear with energies $\omega < 300 - 500$ MeV, i.e. in the very central region of the rapidity space. The number of photons per collision and their spectrum are — apart from logarithms — independent of the

c.m. energy \sqrt{s} . The dependence of the cross section on the photon energy and on the emission angle θ^* with respect to the beam axis is estimated roughly as follows:

$$\frac{1}{\sigma} \frac{d\sigma}{d\omega d\cos\theta^*} \propto \frac{1}{\omega^{1/3}(\sin\theta^*)^{2/3}} \quad (4.1)$$

This should be compared to the inner-bremsstrahlung spectrum

$$\left. \frac{1}{\sigma} \frac{d\sigma}{d\omega d\cos\theta^*} \right|_{\text{bremsstr.}} \propto \frac{1}{\omega \sin^2\theta^*} \quad (4.2)$$

The “synchrotron” radiation from the quarks should thus be harder than the bremsstrahlung spectrum. This is welcome, since for $\omega \rightarrow 0$ bremsstrahlung should dominate according to Low’s theorem [45]. It is amusing to note that in several experiments an excess of soft prompt photons over the bremsstrahlung calculation has been observed [46]–[50]. The gross features and the order of magnitude of this signal make it a candidate for our “synchrotron” process. A detailed comparison with our formulae has been made recently at least for the results from one experiment [50] with encouraging results [51].

One might think — maybe rightly — that these ideas are a little crazy. But we have also worked out some consequences of them for the Drell-Yan reaction (1.2), which make us optimistic. In the lowest order parton process contributing there, we have a quark-antiquark annihilation giving a virtual photon γ^* , which decays then into a lepton pair (Fig. 1):

$$q + \bar{q} \rightarrow \gamma^* \rightarrow \ell^+ \ell^-. \quad (4.3)$$

In the usual theoretical framework q and \bar{q} are assumed to be uncorrelated and unpolarized in spin and colour if the original hadrons are unpolarized. From our ideas we would expect a transverse spin correlation due to the “storage ring” effect [43], [44]. We worked this out and found that this influences the $\ell^+ \ell^-$ angular distribution in a profound way. Then our colleague H. J. Pirner pointed out to us that data which may be relevant in this connection existed already [52]. And very obligingly these data seem to support the idea of spin correlations and thus vacuum effects in high energy collisions. For more details we refer to [53]. If such spin correlations are confirmed by experiments at higher energies, we would presumably have to reconsider the fundamental factorization hypothesis for hard reactions which we sketched in Sect. 1.

5 Conclusions

We have given a sketch of methods and results concerning possible nonperturbative effects in high energy hadronic collisions. I think it is rather sure that such effects

play a decisive role in soft hadronic collisions. I find it even more exciting that they may also influence hard reactions like the Drell-Yan process. In any case it is a challenge for theorists to study such nonperturbative effects in quantum field theory. And here is my outlook and the message which I would like to convey: The **THEORY** of confined quarks may lead to

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perturbation theory predicts. **EXPERIMENT**alists please check.

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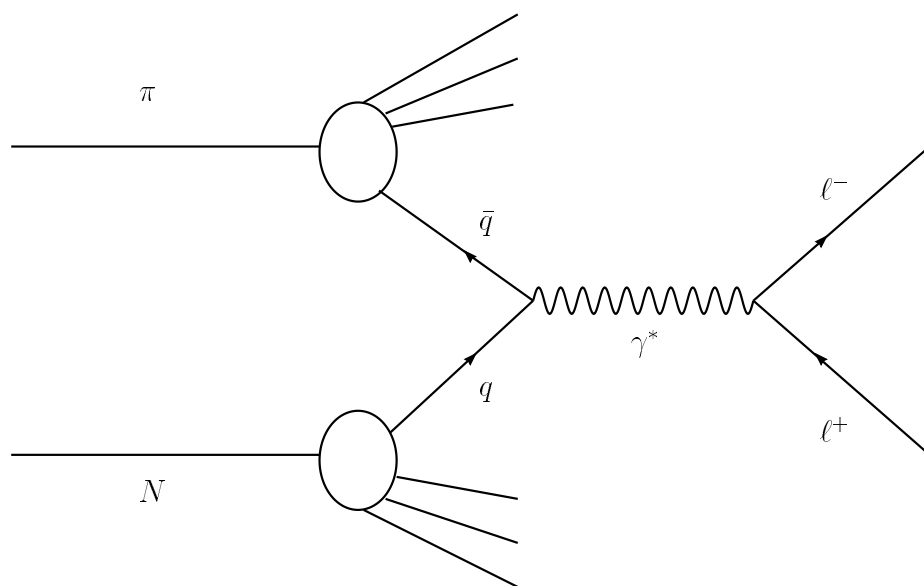
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Figure Captions

- Fig. 1:** The lowest order diagram for the Drell-Yan reaction (1.2) in the QCD improved parton model.
- Fig. 2:** The schematic behaviour of the vacuum energy density $\varepsilon(B)$ as function of a constant chromomagnetic field B according to Savvidy's calculation (eq. (2.2)).
- Fig. 3:** A “snapshot” of the QCD vacuum showing a domain structure of spontaneously created chromomagnetic fields.
- Fig. 4:** The QCD-vacuum according to Ambjørn and Olesen [28] (a). The ether according to Maxwell [33] (b).
- Fig. 5:** An encounter of two quarks: q_1 moving fast in positive and q_2 fast in negative x^3 direction. The interaction of the quarks with the gluons in the QCD vacuum is indicated by the spiral lines.
- Fig. 6:** Proton-proton elastic scattering at high energies with the scattering of single quarks on each other.
- Fig. 7:** Scattering of two quarks q_1, q_2 in a given external gluon potential $G_\lambda(x)$.
- Fig. 8:** Schematic representation of the functional integral on the r.h.s. of (3.4) which sums up the diagrams with arbitrary numbers n and n' of gluon vertices on the quark lines q_1 and q_2 , respectively. This gives the amplitude for $q - q$ scattering.
- Fig. 9:** Simplest approximation for the functional integral on the r.h.s. of (3.4) for $q - q$ scattering without colour exchange.
- Fig. 10:** The relation between the total cross section σ_{tot} and the slope parameter b for proton-proton and proton-antiproton scattering. The dotted line is the prediction from Regge theory. The prediction of the calculation for soft high energy scattering in the stochastic vacuum model is that the data points should lie in the area between the full lines. In essence this is given by (3.8), (3.9) with an uncertainty estimate from different assumptions for the proton wave functions (cf. [41]).
- Fig. 11:** A proton-proton collision at high energies in the parton picture.

Fig. 12: A quark and antiquark traversing a region of chromomagnetic field (a).
An electron and a positron in a storage ring (b). In both cases we expect the emission of synchrotron radiation to occur.

Figure 1



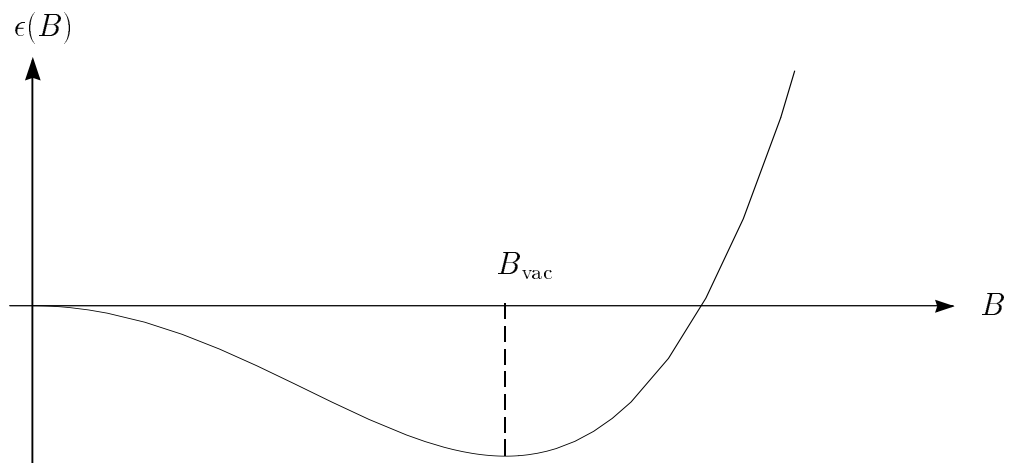
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Figure 2



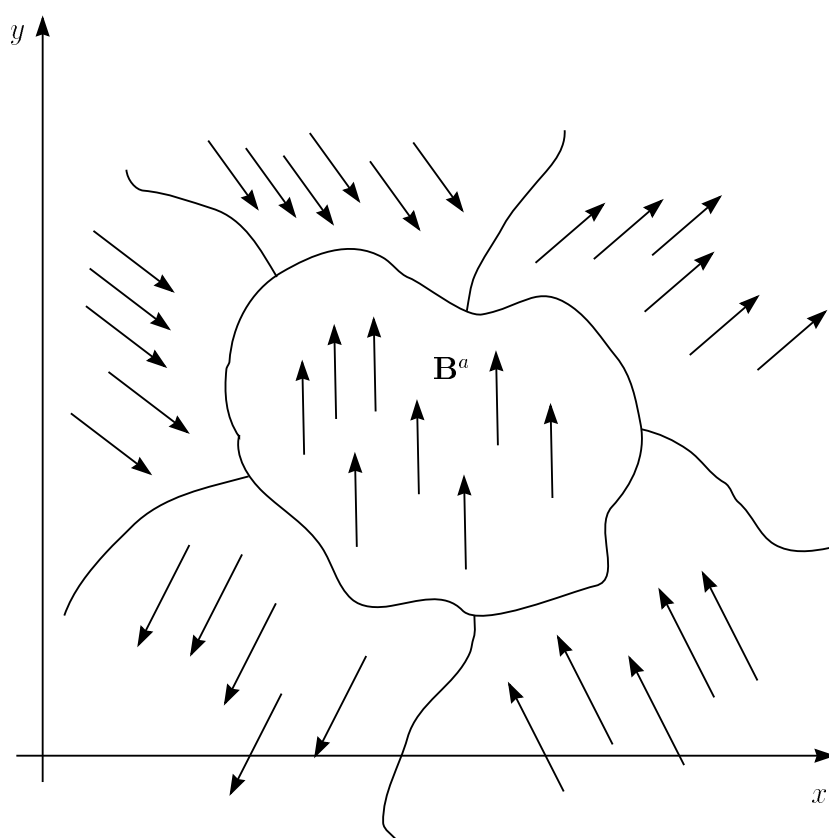
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Figure 3



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Figure 4a

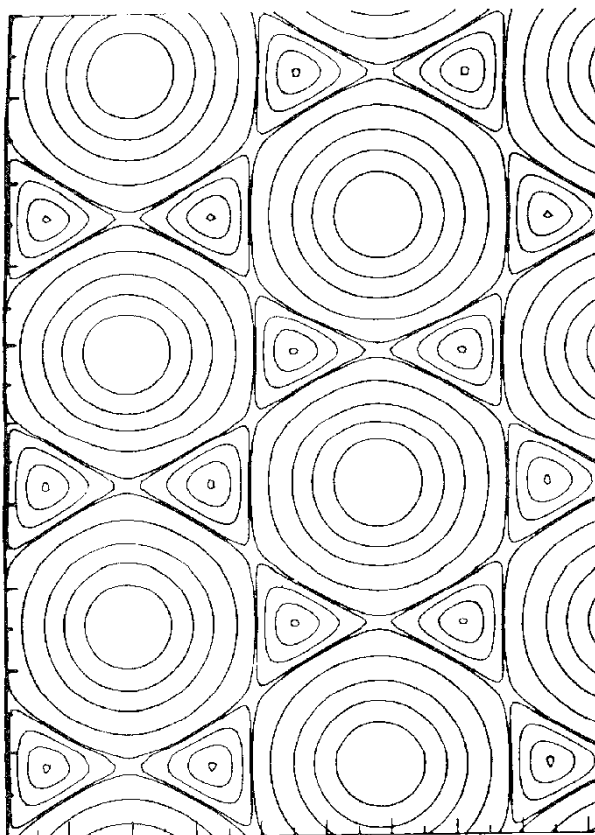


Figure 4b

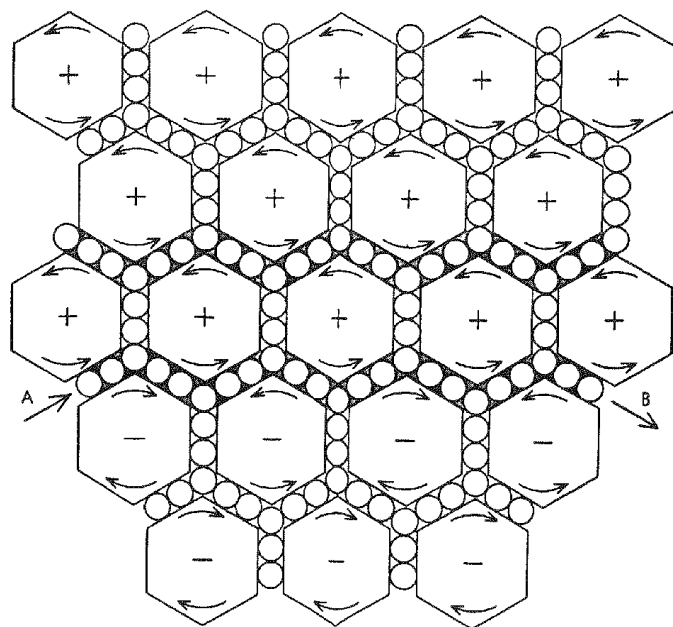


Figure 5

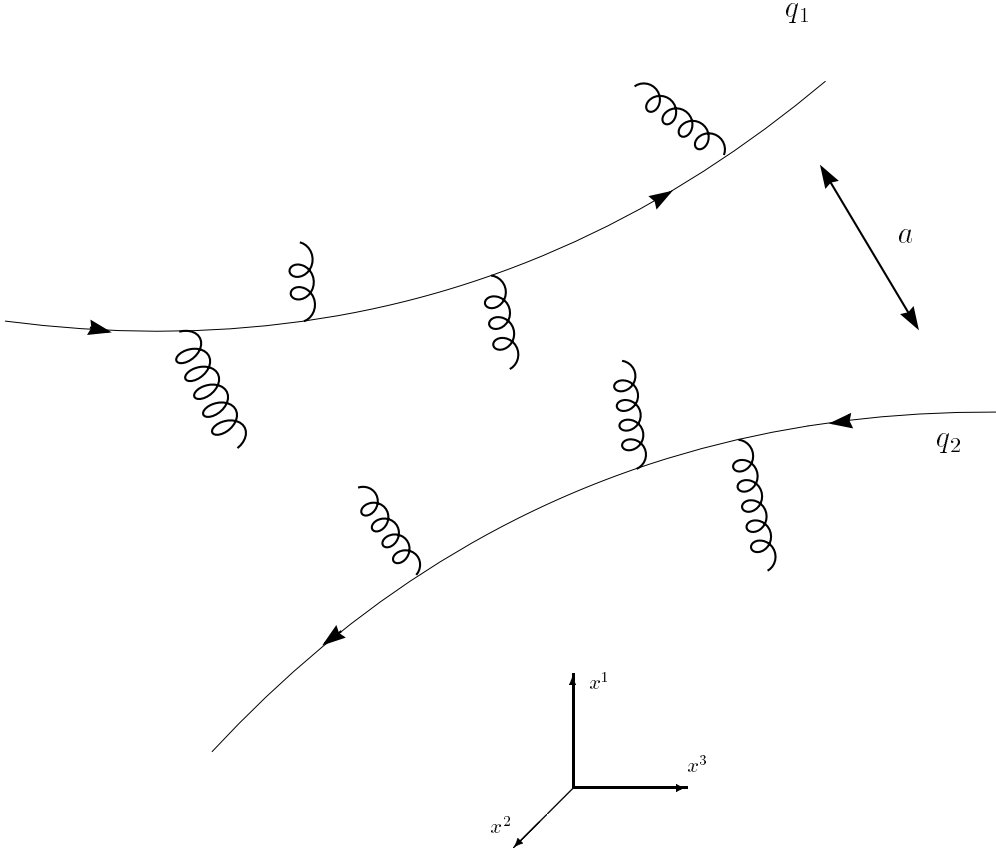


Figure 6

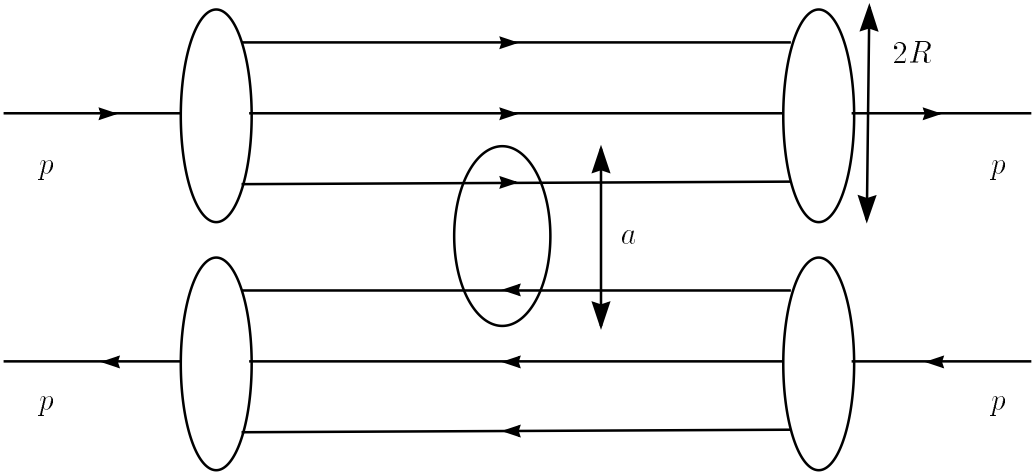


Figure 7

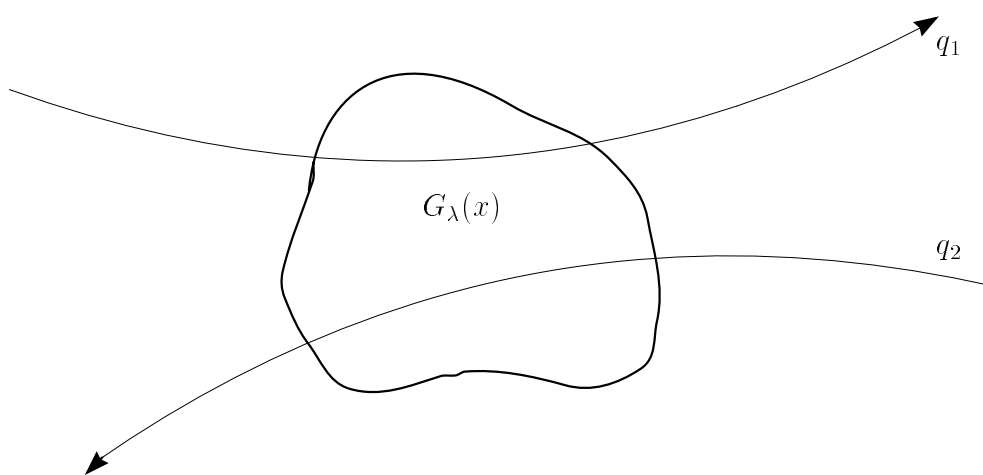


Figure 8

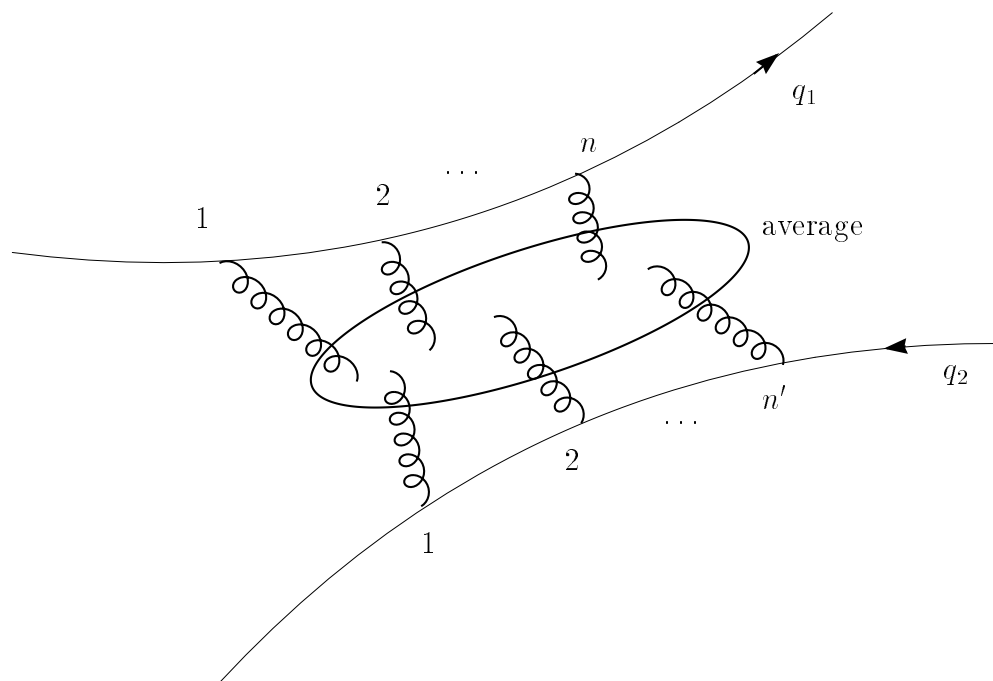


Figure 9

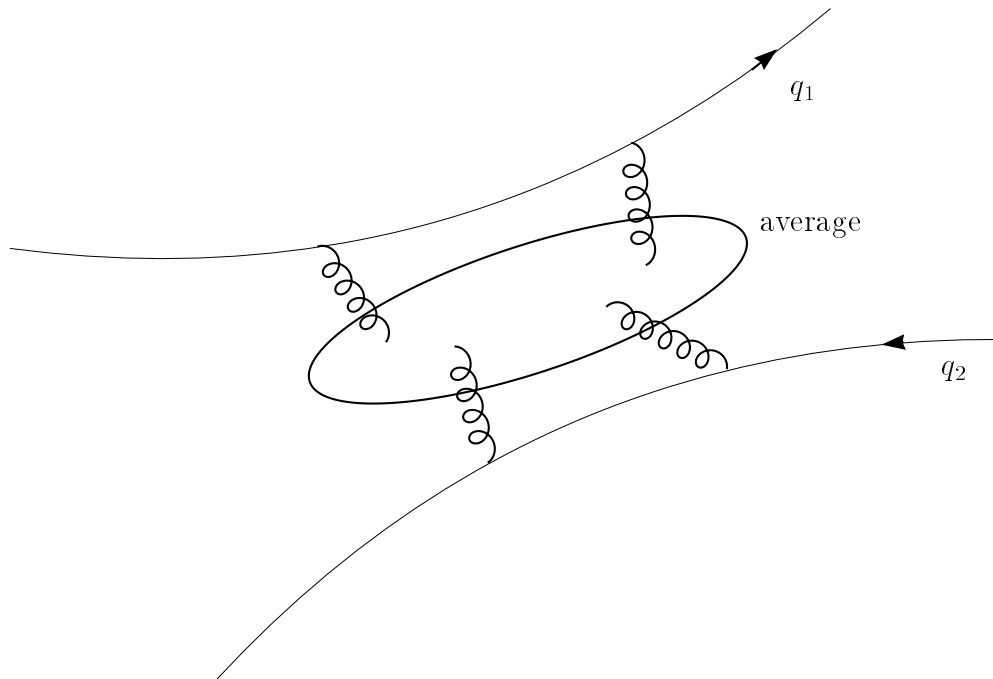


Figure 10

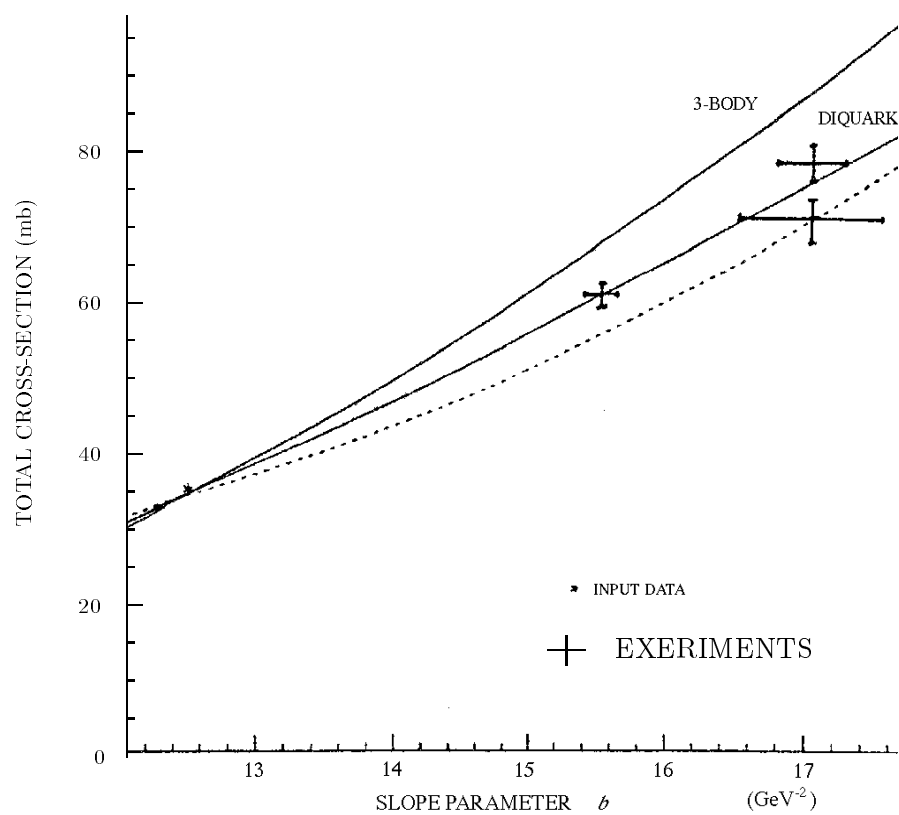


Figure 11

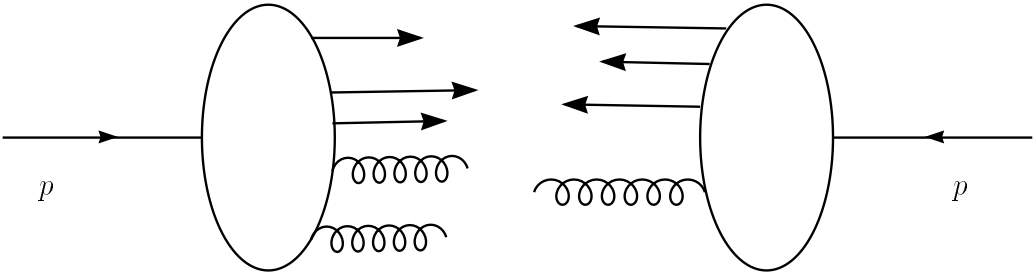


Figure 12

